

Right-handed Sneutrino Inflation in SUSY $B - L$ with Inverse Seesaw

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We propose a scenario for realizing inflation in the framework of supersymmetric $B - L$ extension of the Standard Model. We find that one of the associated right-handed sneutrinos (the superpartner of the Standard Model singlet fermion) can provide a new non-trivial inflationary trajectory at tree level (therefore breaking $B - L$ during inflation). As soon as the inflation ends, the right-handed sneutrino falls into the supersymmetric vacuum, with a vanishing vacuum expectation value, so that $B - L$ symmetry is restored. In this class of models, the $B - L$ gauge symmetry will be radiatively broken at a TeV scale and light neutrino masses are generated through the inverse seesaw mechanism.

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Inflation [1] is one of the most interesting scenarios for explaining the origin of structure formation in the Universe. It leads to a very consistent result with the recent observations on the cosmic microwave background radiation [2]. The basic idea of inflation is based on a scalar field, called inflaton dominating the energy density of the universe, rolling down slowly along an almost flat potential at some early stage of the universe. As a consequence, an extremely rapid exponential expansion of the universe took place. After inflation ends, the field stabilizes in the global (true) minimum of the potential and the universe enters into the radiation dominated era through reheating. Though it is very intriguing to find an inflationary potential from a particle physics model, at the same time this sort of picture is difficult to construct also. The difficulty arises mostly in realizing the flatness of the potential. Several attempts have been made in recent years [3]. There is presently no agreed upon model that can describe inflation from particle physics, however it is now established that the construction requires physics beyond the Standard Model (SM) of particle physics.

An extension of the SM gauge group seems also mandatory for accommodating the non-vanishing neutrino masses confirmed by neutrino oscillation experiments [4]. Recently, it has been shown that the simplest possible such an extension of the SM is the one based on the gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_{B-L}$ [5]. In this class of model, it turns out that the scale of $B - L$ symmetry breaking can be related to the scale of supersymmetry (SUSY) breaking [6]. Therefore, a TeV scale type I or inverse seesaw mechanism can be naturally implemented [7]. In recent years there has been considerable interest in analyzing the phenomenological implications of this class of models and their possible signatures at the Large Hadron Collider (LHC) [8].

In this article we consider the SUSY $B - L$ extension of the SM with inverse seesaw, where right-handed sneutrino plays the role of inflaton. We show that a new non-trivial inflationary trajectory can be realized at tree level, along which the $B - L$ gauge symmetry is broken (dur-

ing inflation), through a large field value of right-handed sneutrino. At the end of inflation, the right-handed sneutrino continues falling into the SUSY vacuum, with a vanishing sneutrino vacuum expectation value (vev). Therefore, $B - L$ symmetry is restored which is important since in the scenario we are interested in (the SUSY $B - L$ extension [7]), the $B - L$ breaking scale is \sim TeV. In our set-up, there is one more field involved in the superpotential for inflation apart from the right handed sneutrino, a singlet under the entire gauge group (X), whose role is to put the inflaton into the right direction during inflation and most importantly to nullify any possible contribution to the neutrino mass matrix by getting a vev at the end of inflation. Our effective inflationary potential then turns out to be similar with the one in chaotic inflation [9] though it is different to start with. Furthermore the dynamics is more involved as we will see later.

Chaotic inflation is a kind of simplest possibility to realize inflation. It can be incorporated in a supersymmetric framework through sneutrino inflation [10]. Here we discuss in brief the basic structure of it which incorporates minimal supersymmetric seesaw model for neutrinos in contrast with our set-up with inverse seesaw mechanism. It is usually assumed that right-handed neutrino superfield, N , is odd under a Z_2 symmetry in order to forbid its cubic term in the superpotential. Therefore the sneutrino effective potential is given by $V = \frac{1}{2} M_N^2 |\tilde{N}|^2$, where M_N is the mass of right-handed neutrino. Such a scenario leads to chaotic inflation, where the universe starts with a random value (chaotic) of the sneutrino inflaton \tilde{N} . It can be noted that in this case the M_N is fixed to be $\sim 10^{13}$ GeV since the amplitude of primordial density perturbation is directly proportional to M_N . In the framework of inverse seesaw, no elements of the neutrino mass matrix is expected to be that large. So we do not expect the above type of sneutrino inflation to take place with inverse seesaw. We have introduced the X field, interplay of which with the right handed neutrino superfield can provide chaotic inflation and simultaneously ensures that there is no such large contri-

butions present in the neutrino mass matrix, thereby the inverse seesaw structure remains intact. With this modification, it also fits the cosmological data. The possibility of hybrid inflation with sneutrino inflaton and with non-minimal Kähler potential has been considered in Ref.[11]. Another class of models where the inflationary direction involves the right handed sneutrino field can be found in [12].

Before going to discuss the inflationary picture, let us discuss in brief the construction of the SUSY $B-L$ extension of the SM with inverse seesaw mechanism. The particle content of this model includes the following superfields in addition to those in the MSSM: two SM singlet chiral Higgs superfields $\chi_{1,2}$ (the vev of their scalar components spontaneously break the $U(1)_{B-L}$ at the TeV scale and χ_2 is required to cancel the $U(1)_{B-L}$ anomaly); three sets of SM singlet chiral superfields N_i, S_{1i}, S_{2i} ($i = 1, 2, 3$) (to implement the inverse see-saw mechanism and also to cancel $B-L$ anomaly). The superpotential of the leptonic sector in this model is given by [7]

$$W = Y_E L H_1 E^c + Y_\nu L H_2 N^c + Y_S N^c \chi_1 S_2 + \mu H_1 H_2 + \mu' \chi_1 \chi_2. \quad (1)$$

The charge specification of superfields, appeared in this superpotential, under $U(1)_{B-L}$ and $U(1)_R$ are given in Table I. A term proportional to $S_1 S_2$ is also allowed by

Charges	E^c	L	H_1	H_2	N^c	χ_1	χ_2	S_1	S_2
$B-L$	1	-1	0	0	1	1	-1	2	-2
R	2	0	0	2	0	1	1	1	1

TABLE I: Charges under $U(1)_{B-L}$ and $U(1)_R$.

the symmetry, which we will discuss later in the context of inflationary superpotential along with the term involved with X field.

As mentioned above, the $B-L$ symmetry is radiatively broken by the non-vanishing VEVs $\langle \chi_1 \rangle = v'_1$ and $\langle \chi_2 \rangle = v'_2$ [6]. After $B-L$ and EW symmetry breaking, the neutrino Yukawa interaction terms lead to the following expression:

$$\mathcal{L}_m^\nu = m_D \bar{\nu}_L N^c + M_N \bar{N}^c S_2 + \text{h.c.}, \quad (2)$$

where $m_D = Y_\nu v \sin \beta$, $M_N = Y_S v' \sin \theta$ and $v'^2 = v_1'^2 + v_2'^2$. In this framework, the light neutrino masses are related to a small mass term $\mu_s S_2^2$ in the Lagrangian, with $\mu_s \sim \mathcal{O}(1)$ KeV which can be generated at $B-L$ scale through non-renormalizable higher order term $\frac{\chi_1^4 S_2^2}{M_I^3}$, with M_I is an intermediate scale of order $\mathcal{O}(10^7)$ GeV. Therefore, the following light and heavy neutrino masses are obtained respectively:

$$m_{\nu_e} = \frac{m_D^2 \mu_s}{M_N^2 + m_D^2}, \quad (3)$$

$$m_{\nu_{H,H'}} = \pm \sqrt{M_N^2 + m_D^2} + \frac{1}{2} \frac{M_N^2 \mu_s}{M_N^2 + m_D^2}. \quad (4)$$

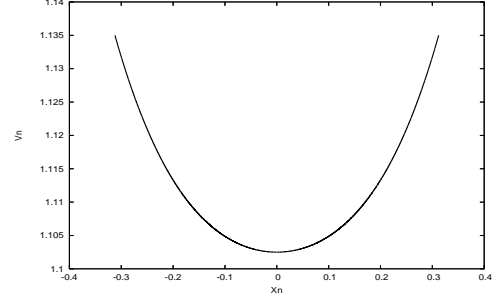


FIG. 1: Shape of the Potential ($V_n = V/(2m^3 M_*)$) along X direction (X scaled as $X_n = X/(m M_*)^{1/2}$) for $\sigma = 1.05\sigma_c$.

A noticeable difference with the usual seesaw mechanism is that in this case, no terms in the neutrino mass matrix need to be very large.

In this type of models, the sneutrino is given by linear combination of $(\tilde{\nu}_L, \tilde{\nu}_L^\dagger, \tilde{N}, \tilde{N}^\dagger, \tilde{S}_2, \tilde{S}_2^\dagger, \tilde{S}_1, \tilde{S}_1^\dagger)^T$. In general, the \tilde{S}_1 's are decoupled and have no interactions with the SM particles. In Ref.[13], it was shown that the lightest right-handed sneutrino in this model is a viable candidate for cold dark matter. Also, it was emphasized [14] that the one-loop radiative corrections due to right-handed (s)neutrinos to the mass of the lightest Higgs boson is significantly large and can push the upper bound imposed on its mass to around 200 GeV.

Here we show that right-handed sneutrinos $S_{1,2}$ can be the inflaton fields and provide a mechanism for chaotic inflation. For this propose, a gauge singlet superfield, X , is introduced (can be a part of a hidden sector). The superpotential part, which is relevant for inflation, is given by

$$\delta W = S_1 S_2 \left(\frac{X^2}{M_*} - m \right), \quad (5)$$

where m is a mass parameter (can be large enough) in the model which would be responsible for providing the vacuum energy during inflation and M_* is a superheavy scale (can be thought of as some kind of messenger mass between the hidden and $B-L$ sector). Note that the renormalizable coupling $X S_1 S_2$ is not allowed by \mathbb{Z}_2 symmetry, under which X is odd and all other superfields are even. X has R charge zero. Such a renormalizable coupling between the $B-L$ sector and singlet X superfield would not lead to a successful model of inflation. As we have discussed before, our aim is not to break the $B-L$ symmetry at the end of inflation, since $B-L$ needs to be broken at a much lower scale \sim TeV.

The corresponding scalar potential can be written as

$$V = \left| \frac{X^2}{M_*} - m \right|^2 (|S_1|^2 + |S_2|^2) + 4|S_1 S_2|^2 \frac{|X|^2}{M_*^2} + \text{D-terms}, \quad (6)$$

where the scalar component of the superfields are denoted by the same symbols as the corresponding superfields. Note that choosing the direction $|S_1| = |S_2|$ would

lead to vanishing of the D-terms relevant for the discussion of the inflation sector fields. Specifying this direction as $S = |S_1| = |S_2|$ and choosing the real part of it as $S = \sigma/\sqrt{2}$, the potential takes the form,

$$V = \left(\frac{X^2}{M_*} - m \right)^2 \sigma^2 + \sigma^4 \frac{X^2}{M_*^2}, \quad (7)$$

where the phase of X is taken to be zero for simplicity. For nonzero σ , the potential above have extrema along,

$$X = 0 \quad \text{or} \quad 2X^2 + \sigma^2 = 2mM_*, \quad (8)$$

directions. We note that for $\sigma > \sigma_c = (2mM_*)^{1/2}$, $X = 0$ corresponds to a local minimum. At this minimum, the potential V takes the form $V_{\text{inf}} = m^2\sigma^2$, which can provide inflation (σ is our inflaton). The change in the shape of the potential (obtained from δW) along X direction can be understood from Fig.1 and 2 for different values of S , above and below S_c respectively.

It is interesting to note that V_{inf} is similar to the potential responsible for chaotic inflationary picture and therefore the inflation can be realized for large value of σ above the reduced Planck scale: $M_P = 2.4 \times 10^{18}$ GeV. However in our case, the dynamics of having this potential is different compared to the chaotic inflation and there are more fields involved here. As we have mentioned at the beginning, the construction of the inverse seesaw mechanism should not be destabilized by the additional large mass parameter m introduced in δW ($m \sim 10^{13}$ GeV as we will find soon). This can indeed be achieved if X can acquire a vev ($\langle X \rangle$) at the end of inflation. As the expected contribution to the neutrino mass matrix from δW is proportional to $(X^2 - mM_*)$, a vev of X , $\langle X \rangle = (mM_*)^{1/2}$, at the end of inflation would ensure the stability of the neutrino mass through the inverse seesaw. As we proceed further, we will find how X can acquire such a vev at the end of inflation. So the light neutrino mass is essentially independent of the mass parameter involved in the inflationary potential as opposed to the case with sneutrino inflation [10]. This also explains why a single term in the superpotential mS_1S_2 which can provide a successful chaotic sneutrino inflation as can not be accommodated alone while working with inverse seesaw mechanism.

At the end of inflation, the field σ rolls down towards the global minimum at $\sigma = 0$. As soon as σ is below σ_c , $X = 0$ becomes a local maximum and a new valley for the X field appears along $X^2 = mM_* - \sigma^2/2$ direction, which would be the new local minima. Therefore X is trapped along that direction only and will stay at this valley by continuously adjusting itself in this direction with the change of σ . It is therefore expected to make a smooth transition to $\langle X \rangle = (mM_*)^{1/2}$ at the very end when σ becomes very small and subsequently goes to zero. However we notice that the global minimum lies at $\sigma = 0$, thereby making X arbitrary. This arbitrariness is actually corrected once we assume a contribution from a hidden sector potential carrying interaction (suppressed though) with the X field. We will come back on this

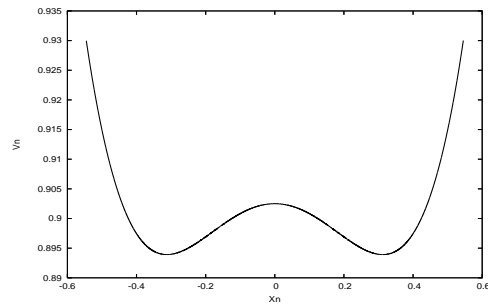


FIG. 2: Shape of the Potential ($V_n = V/(2m^3M_*)$) along X direction (X scaled as $X_n = X/(mM_*)^{1/2}$) for $\sigma = 0.95\sigma_c$.

issue while discussing the post inflationary phase near the end of this article. With this new contribution appearing from hidden sector (not playing any role in inflation), $\langle X \rangle$ will remain same as $(mM_*)^{1/2}$, which is instrumental for keeping light neutrino mass matrix uninterrupted.

We can now discuss the inflationary parameters [3] in connection with experimental observations. During inflation, X is at zero and the potential is dominated by $V_{\text{inf}} = m^2\sigma^2$. The slow roll parameters are given by

$$\epsilon = \frac{m_{Pl}^2}{16\pi} \left(\frac{V'_{\text{inf}}}{V_{\text{inf}}} \right)^2 \approx \frac{m_{Pl}^2}{4\pi\sigma^2}, \quad \eta = \frac{m_{Pl}^2}{8\pi} \left(\frac{V''_{\text{inf}}}{V_{\text{inf}}} \right) \approx \frac{m_{Pl}^2}{4\pi\sigma^2}, \quad (9)$$

where m_{Pl} is the Planck scale $\sim 1.22 \times 10^{19}$ GeV. Note that the slow roll parameters are independent of the mass scale m . The number of e-foldings is provided through the relation,

$$N = \frac{8\pi}{m_{Pl}^2} \int_{\sigma_e}^{\sigma_0} \frac{V_{\text{inf}} d\sigma}{V'_{\text{inf}}} \approx \frac{2\pi}{m_{Pl}^2} (\sigma_0^2 - \sigma_e^2), \quad (10)$$

where the subscript e corresponds to the values at the end of inflation and 0 indicates the values when comoving scales comparable to our present horizon crossed outside the inflationary horizon (*i.e.* when the comoving wavenumber $k_0 = 0.002$ Mpc $^{-1}$). Hence the slow roll parameter at that time can be expressed as $\epsilon_0 = 1/(1 + 2N)$, where we have used the fact that $\epsilon_e \sim 1$. So with $N \simeq 60$, $\epsilon_0 \simeq 0.008$ at that time. The amplitude of curvature perturbation Δ_R is given by,

$$\Delta_R = \frac{8}{m_{Pl}^3} \sqrt{\frac{2\pi}{3}} \frac{V_{\text{inf}}^{3/2}}{|V'_{\text{inf}}|}. \quad (11)$$

According to WMAP seven years data[2], $\Delta_R \sim 4.93 \times 10^{-5}$ corresponding to k_0 . Therefore from eq.(10) and eq.(11), we find that $m \sim 1.3 \times 10^{13}$ GeV. Now using the expression for the spectral index for scalar density perturbation, $n_s = 1 - 6\epsilon + 2\eta$, we find $n_s \sim 0.97$, which is in good agreement with recent WMAP data [2]. The prediction for the ratio of scalar and tensor perturbations (r) at the quadrupole scale is found to be $r \sim 0.13$ which is consistent with the WMAP findings [2].

We are now in a position to discuss the post inflationary phase and reheating. We have seen that for $\sigma < \sigma_c$, X is trapped in the valley, $X^2 = mM_* - \sigma^2/2$ and when σ approaches zero, the field X should smoothly enter into the minimum at $X = (mM_*)^{1/2}$. However $S = 0$ (or $\sigma = 0$) makes X arbitrary as at this stage X field is massless. It results to a possibility that X can be deviated from its would be minimum at $(mM_*)^{1/2}$, which is unwanted for keeping the neutrino mass in the right range. In order to cure this problem, we introduce a term in the scalar potential, originated from an interplay between a hidden sector and X ,

$$V_h(X, Y) = h^2 |Y|^2 (X - M)^2, \quad (12)$$

where Y is a superfield in the hidden sector, M is a mass scale related by $M^2 = mM_*$ and h is the coupling constant. This choice of M ensures that X minimum is not shifted anymore from $X = (mM_*)^{1/2}$ provided $\langle Y \rangle \neq 0$. This relation may not be completely unnatural as the other mass parameter involved in the scenario, M_* , is not restricted from inflation. At this point we do not attempt to construct any specific model for the hidden sector so as to realize V_h directly. We assume that Y get this vev when $\sigma < \sigma_c$, so that the inflationary phase is not perturbed by the introduction of V_h (during inflation Y is expected to be at the origin).

Therefore in this picture, during early stage of inflation and thereafter, both S and X have field dependent masses which are decreasing as S approaches zero. Following our assumption, at some stage (when X is already in the valley of $X^2 = mM_* - \sigma^2/2$), X starts to receive an additional contribution towards its mass from V_h with $h\langle Y \rangle = m_X$. This massive X will oscillate around its minimum and can therefore decay and reheat the universe. From eq.(5), we note that a decay of X to fermionic components of S_1, S_2 is possible and can be estimated as $\Gamma_1 \approx (h/\pi)(M/M_*)^2 \langle Y \rangle$. For a choice of $M_* \sim 10^{16}$ GeV (close to the GUT scale) with $h \sim O(10^{-2})$, the reheat temperature, $T_R \simeq \frac{1}{7}\sqrt{\Gamma M_P}$, is found to be $\lesssim 10^9$ GeV provided $\langle Y \rangle \lesssim 10^7$ GeV.

From eq.(12), it can be found that the X field may also decay into Y and thereby reheating could be into the hidden sector only, provided this decay mode is the most efficient. However in order to achieve a successful reheating of the observable universe, we can assume that the mass of the scalar component of the Y field (m_Y) is bigger than m_X . So this decay and hence the reheat-

ing into the hidden sector only can be prevented. We can safely assume this ($m_Y > m_X$) as m_Y would depend on the constructional details of the hidden sector itself. There could be another way out if the Y field is part of the supersymmetry breaking sector. In this case through the presence of the mediator of the SUSY breaking connecting the hidden as well as the observable sector, the universe can be finally reheated as shown in [15]. Another possibility is the preheating [16] of the universe after inflation through the parametric resonance. However this would involve several other fields in our setup and requires detailed study. At this point, the detailed discussion of reheating is beyond the scope of this article.

In conclusion, we have constructed a model for sneutrino inflation embedded in SUSY $B-L$ extension of the SM along with an additional singlet field, where the light neutrino mass is generated through inverse seesaw mechanism. At an early stage of the universe, these sneutrino fields are having very large values and inflation is initiated. $B-L$ symmetry is therefore broken at that time. However when inflation is over, the inflaton rolls down to zero and additional singlet field receives a nonzero vacuum expectation value at the very end, thereby the $B-L$ symmetry is restored. This $B-L$ symmetry will be broken radiatively at a TeV scale and the electroweak symmetry breaking will take place subsequently. Then the light neutrino mass is generated at the expected level. It can be noted that the interplay of two fields involved in realizing inflation bears a similarity with the supersymmetric hybrid inflation models (SHI) [17]. In SHI, the superpotential is linear in the inflaton superfield. So the inflationary direction is absolutely flat and a slope needs to be generated using the Coleman-Weinberg correction. On the contrary, there is a nonvanishing potential existing for the inflaton in our case along which it rolls down naturally towards its minimum at zero.

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